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Higgs-dependent Leptogenesis

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We present and study a novel electroweak-scale leptogenesis mechanism occurring in $U(1)_F$ flavor symmetric two Higgs doublet models with Higgs-dependent Yukawa couplings. In this scenario CP-violation originates entirely from the Higgs sector, and large CP asymmetries in TeV scale heavy neutrino decays can be obtained without resonance enhancement and any fine tuning of model parameters. Distinctive predictions for fermion Yukawa couplings together with CP-violating Higgs sector make tests of the electroweak-scale leptogenesis mechanism realistic at LHC experiments.

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INTRODUCTION

The origin of quark and lepton Yukawa couplings remains the major unsolved question in particle physics [1]. One commonly accepted approach, the Froggatt-Nielsen mechanism [2], explains the observed fermion mass hierarchies with powers of flavour symmetry $U(1)_F$ breaking vacuum expectation value (VEV), $(\langle S_F\rangle/\Lambda_F)^n$, where Λ_F is the scale of new physics, and n is determined by fermion quantum numbers under $U(1)_F$. An interesting recent proposal [3, 4] suggests that all Yukawa couplings are generated by the SM Higgs boson H via the operators $c_n(H^\dagger H/M^2)^n$. This scenario *predicts* the flavor physics scale $M \sim \mathcal{O}(1)$ TeV. However, as $H^\dagger H$ is a $U(1)_F$ singlet, no Froggatt-Nielsen type models can be built with one H, and the Yukawa hierarchies are based on the additional assumption that some coefficients c_n vanish or on some unknown non-Abelian symmetry [4].

Leptogenesis [5] is the most promising mechanism to explain the observed [6] baryon asymmetry of the Universe. In the standard scenario [5] the CP asymmetry ε_{N_i} in heavy Majorana neutrino N_i decays is induced by complex neutrino Yukawa couplings y^ν_{ij}. In the case of $\mathcal{O}(1)$ TeV scale heavy neutrinos the seesaw mechanism [7] implies $y^\nu \lesssim 10^{-7}$, and large enough CP asymmetry cannot be generated unless masses of the heavy neutrinos satisfy the resonance condition [8], which requires unnatural fine tuning between M_{N_i} at low scale.

In this Letter we propose a model for Higgs-dependent Yukawa couplings which successfully explains the masses and mixings of quarks, charged leptons and neutrinos. We extend the SM Higgs sector with another doublet. Because H_uH_d is nontrivial under $U(1)_F$, Froggatt-Nielsen type model building defines Yukawa coupling hierarchies from fermion flavor quantum numbers [4]. In two Higgs doublet models the CP-violation may come from the Higgs sector [9]. We show that in this scenario new leptogenesis mechanism occurs below the scale of electroweak symmetry breaking (EWSB), $T < T_c$, in which the necessary CP violation comes entirely from the Higgs potential. The CP asymmetry is *not* suppressed by small neutrino Yukawa couplings and successful leptogenesis occurs *naturally* for non-degenerate heavy neutrino masses *independently* of the initial conditions on N_i abundances. Taking into account flavor effects [10] in all heavy neutrino decays we derive and solve Boltzmann equations for the lepton asymmetry as well as for sphaleron transitions [11]. Although sphalerons decay soon below the EWSB scale, $T_d < T_c$ [12, 13], we show that the lepton asymmetry is safely converted to the baryon asymmetry. This scenario implies a phenomenological *upper bound* on M_{N_i} from successful leptogenesis. Because the Higgs sector CP-violation is testable at the LHC [9], our proposal opens completely new, *realistic,* perspectives for testing electroweak-scale leptogenesis mechanisms at collider experiments.

THE MODEL

We consider global $U(1)_F$ symmetric two Higgs doublet model with flavor charges given by $Q(H_u) = h_u$, $Q(H_d) = h_d$, $Q(L_i) = -4h_u - 3h_d$ and $Q(N_{Ri}) = 0$. We assume $h_u + h_d \neq 0$ so that the term H_uH_d is not $U(1)_F$ invariant. The $U(1)_F$ invariant Yukawa Lagrangian for heavy Majorana neutrinos N_{Ri} is given by

$$\mathcal{L} = y^\nu_{ij}\bar{N}^i_R L^j_L H_u\left(\frac{H_uH_d}{M^2}\right)^{n^\nu_{ij}} - \frac{1}{2}N^i_R M^{ij}_N N^j_R + h.c., \quad (1)$$

and similarly for quarks and charged leptons. In Eq. (1) we have neglected the term $\bar{N}^i_R L^j_L H^\dagger_d\left(H_uH_d/M^2\right)^{n^\nu_{ij}+1}$ because its contribution to the induced Yukawa couplings is additionally suppressed. Eq. (1) gives rise to $(2n^\nu+1)$ Higgs boson interactions with N and L. Below the EWSB scale the Higgs bosons acquire VEVs $\langle H_u\rangle = v_u = v\sin\beta$ and $\langle H_d\rangle = v_d = v\cos\beta$, where $v = 174$GeV. The small flavor symmetry violating parameter in this scenario is $\epsilon \equiv v_uv_d/M^2 \sim 10^{-2}$, and for $\tan\beta = 1$ the cut-off scale of this model is $M = 1.23$ TeV. Expanding the neutral components of Higgs fields $(v+H_0)^{2n^\nu+1}/M^{2n^\nu}$ in Eq. (1), we get the fermion mass terms proportional to $\epsilon^{n^\nu}v$ and multi-Higgs interaction terms $\epsilon^{n^\nu}H_0(H_0/v)^k$, $k = 0,..,2n^\nu$, which have the common suppression factor ϵ^{n^ν}. The existence of quadratic interaction $\epsilon^{n^\nu}H_0^2N\nu/v$ below EWSB plays a crucial role in our scenario.

Under the charge assignment presented above, n^ν_{ij} are universally given by $n^\nu_{ij} = 3$. This gives the correct

pattern for neutrino Dirac Yukawa couplings and, together with the Majorana mass terms, imply the seesaw induced light neutrino mass scale $m_\nu \sim v^2 s_\beta^2 \epsilon^6/M_N = 1.5\times 10^{-2}$ eV for $M_N = 1$TeV. As is common to Abelian flavor models, the observed neutrino masses and mixing are obtained by appropriate adjustment of the $\mathcal{O}(1)$ coefficients $y_{ij}^{\nu,e}$. In the charged lepton sector the flavor charges are given by $n_{1j}^e = 3$, $n_{2j}^e = 2$ and $n_{3j}^e = 1$, which imply the measured masses via $m_\tau \sim \epsilon v c_\beta$, $m_\mu \sim 6\epsilon^2 v c_\beta$ and $m_e \sim 3\epsilon^3 v c_\beta$. Similarly, the correct quark masses and mixing are obtained with flavor charges $n_{i1}^d = n_{i2}^d = 2$, $n_{i3}^d = 1$, $n_{1j}^u = (2,2,1)$, $n_{2j}^u = n_{3j}^u = (1,1,0)$.

The Higgs potential of this model is given by

$$\begin{aligned} V &= m_{H_u}^2 |H_u|^2 + m_{H_d}^2 |H_d|^2 + \lambda_1 |H_u|^4 + \lambda_2 |H_d|^4 \\ &+ \lambda_3 |H_u|^2 |H_d|^2 + \lambda_4 |H_u H_d|^2 \qquad (2) \\ &+ \Big[m^2 H_u H_d + \lambda_5 (H_u H_d)^2 + \lambda_6 |H_u|^2 H_u H_d \\ &+ \lambda_7 |H_d|^2 H_u H_d + h.c. \Big], \qquad (3) \end{aligned}$$

where Eq. (2) is $U(1)_F$ symmetric and Eq. (3) is explicitly $U(1)_F$ breaking part of V. While $m_{H_{u,d}}^2$ and $\lambda_{1,2,3,4}$ are real, m^2 and $\lambda_{5,6,7}$ are in general complex and give rise to CP-violation. Minimization of the potential V requires

$$b \equiv m^2 - v^2 s_\beta c_\beta (\lambda_4 + 2\lambda_5 - \lambda_6 \tan\beta - \lambda_7 \cot\beta) \qquad (4)$$

to be real and the vacuum stability condition implies

$$\begin{gathered} \lambda_1 > 0,\ \lambda_2 > 0,\ \lambda_3 + \sqrt{\lambda_1\lambda_2} > 0, \\ \lambda_3 + \lambda_4 + 2\sqrt{\lambda_1\lambda_2} - 2|\lambda_5| > 0 \text{ (for } \lambda_6 = \lambda_7 = 0). \quad (5) \end{gathered}$$

The mass matrix $(1/2) H_0^T M_0^2 H_0$ for neutral Higgs bosons $H_0 = (\mathrm{Re}H_u^0, \mathrm{Re}H_d^0, \mathrm{Im}H_u^0, \mathrm{Im}H_d^0)^T$ has one zero eigenvalue which corresponds to the Goldstone boson eaten by Z. M_0^2 is diagonalized by the orthogonal transformation $O^T M_0^2 O = \mathrm{diag}(m_{h_1}^2, m_{h_2}^2, m_{h_3}^2, 0)$, and the mass eigenstates h_a, $a = 1,..,4$, are obtained from the weak eigenstates H_0 via $h = O^T H_0$. Since CP is violated in the Higgs potential by complex parameters, real and imaginary components of the neutral Higgs bosons are mixed with each other. If all couplings were real, off-diagonal blocks of M_0^2 vanish, and $h_{1,2}$ and h_3 would correspond to the CP even and odd Higgs bosons, respectively.

Inserting Higgs mass eigenstates into Eq. (1) implies $N_R - \nu_L$ interactions up to dimension-five operators as

$$\mathcal{L}_\nu = \bar{N}^i P_L (U_{MNS}\,\nu)^j \left(A_{ij}^a h_a + \frac{1}{v} B_{ij}^{ab} h_a h_b \right) + h.c., \qquad (6)$$

where the effective Yukawa couplings A and B are

$$\begin{aligned} A_{ij}^a &= \frac{(-1)^{n_{ij}^\nu}}{\sqrt{2}} y_{ij}^\nu \epsilon^{n_{ij}^\nu} \left[(1 - n_{ij}^\nu)(O_{1a} + iO_{3a}) \right. \\ &- \left. n_{ij}^\nu (O_{2a} + iO_{4a}) \right], \qquad (7) \\ B_{ij}^{ab} &= \frac{(-1)^{n_{ij}^\nu}}{8} y_{ij}^\nu \epsilon^{n_{ij}^\nu} \left[\frac{(O_{1a} + iO_{3a})}{s_\beta} + \frac{(O_{2a} + iO_{4a})}{c_\beta} \right] \\ &\times \left[n_{ij}^\nu (n_{ij}^\nu - 3)(O_{1b} + iO_{3b}) \right. \\ &+ \left. n_{ij}^\nu (n_{ij}^\nu - 1)(O_{2b} + iO_{4b}) c_\beta \right] + (a \leftrightarrow b). \qquad (8) \end{aligned}$$

The couplings A and B are complex because O as well as y^ν are complex. However, the CP asymmetry in heavy neutrino decays depends on $|y^\nu|^2$ and, therefore, only the phases in O contribute to the CP asymmetry.

The $N_R - e_L$ interactions are given by

$$\mathcal{L}_e = (-1)^{n_{ij}^\nu} \epsilon^{n_{ij}^\nu} y_{ij}^\nu \bar{N}_i P_L e_j \phi^+ + h.c., \qquad (9)$$

where ϕ^+ is the charged Higgs boson with mass $2b/\sin 2\beta$. The three-point vertex for neutral Higgs bosons is

$$V_3 = v\ C_{abc} h_a h_b h_c, \qquad (10)$$

where the CP-violating couplings C_{abc} are lengthy expressions of the parameters in V.

LEPTOGENESIS

In our scenario the heavy neutrino decays $N \to LH$ are induced below EWSB. The usual leptogenesis CP asymmetry induced by N_i loops [5] is negligible for non-degenerate $\mathcal{O}(1)$ TeV neutrinos. However, in neutral decay modes $N_i \to \sum_{a=1}^3 \nu_j h_a$ new CP asymmetry

$$\varepsilon_i^j = \frac{\sum_a \left[\Gamma(N_i \to \nu_j h_a) - \Gamma(N_i \to \bar{\nu}_j \bar{h}_a) \right]}{\sum_{k,b} \left[\Gamma(N_i \to \nu_k h_b) + \Gamma(N_i \to \bar{\nu}_k \bar{h}_b) \right]} \equiv \frac{N_i^j}{\sum \Gamma_{Di}^{\nu k}} \quad (11)$$

is generated by interactions (6) and (10) via the diagrams in FIG. 1. In Eq. (11)

$$\begin{aligned} N_i^j &= -\frac{1}{16\pi^2} \frac{1}{16\pi} M_{Ni} \sum_{a,b,c=1}^{3} C_{abc} \mathrm{Im}\left[J_{a,bc} \right] \\ &\times \mathrm{Im}\left[(A^a U_{MNS})_{ij} \left(B^{bc} U_{MNS} \right)_{ij}^* \right] \left(1 - \frac{m_{h_a}^2}{M_{Ni}^2} \right)^2, \quad (12) \\ \Gamma_{Di}^{\nu k} &= \frac{1}{16\pi} M_{Ni} \sum_{a=1}^{3} |(A^a U_{MNS})_{ik}|^2 \left(1 - \frac{m_{h_a}^2}{M_{Ni}^2} \right)^2, \quad (13) \end{aligned}$$

where $J_{a,bc}$ is the loop function defined by

$$J_{a,bc} = \int_0^1 dx \ln\left[x(x-1) + (1-x) r_b + x r_c \right], \qquad (14)$$

where $r_{b,c} = m_{h_{b,c}}^2 / m_{h_a}^2$. The function (14) has absorptive imaginary part when the mass of final state h_a is larger than the sum of intermediate Higgs masses, $m_{h_b} + m_{h_c} < m_{h_a}$. These states induce the CP asymmetry (11).

To exemplify the magnitude of generated CP asymmetry we assume non-degenerate heavy neutrino masses $(M_N)_{ij} = M_N \mathrm{diag}(0.5, 1.25, 1.5)$, $M_N = 1$ TeV, and

$$\begin{aligned} \lambda_1 &= 0.2,\ \lambda_2 = 0.5,\ \mathrm{Im}\lambda_5 \neq 0, \\ m^2 &= (300\mathrm{GeV})^2 + 2iv^2 s_\beta c_\beta \mathrm{Im}\lambda_5,\ \text{others} = 0, \quad (15) \end{aligned}$$

and $\tan\beta = 1$. The imaginary part of m^2 is constrained by the minimization condition Eq. (4). For this choice the CP violating coupling $\mathrm{Im}\lambda_5$ is the only free parameter. The neutral Higgs boson masses are approximately

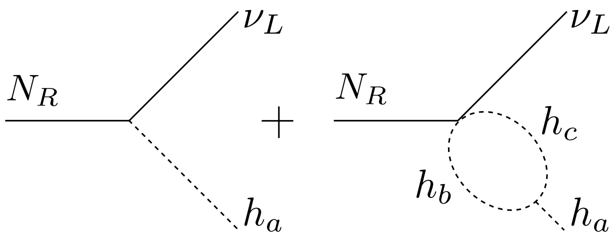

FIG. 1: Diagrams generating dominant contribution to CP asymmetry in heavy Majorana neutrino decay below EWSB.

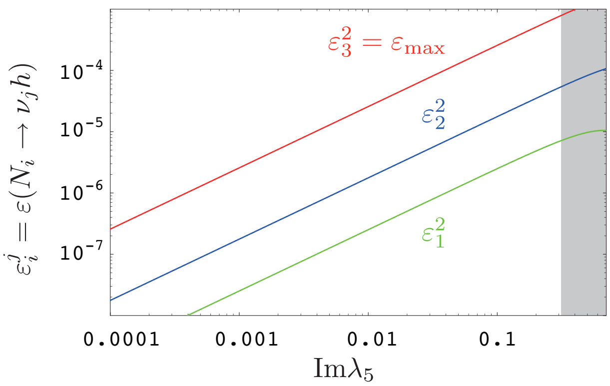

FIG. 2: Examples of CP asymmetries ε_i^j as functions of $\mathrm{Im}\lambda_5$. Decays $N_3 \to \sum_a \nu_2 h_a$ give the largest contribution to ε.

$h_1 \sim 140\mathrm{GeV}$, $h_2 \sim 480\mathrm{GeV}$, $h_3 \sim 440\mathrm{GeV}$ and depend weakly on the magnitude of $\mathrm{Im}\lambda_5$. Therefore, the N_i decay modes to $h_{2,3}$ (with h_1 in the loop) have the absorptive part and non-vanishing CP asymmetry.

Three components of nine CP asymmetries ε_i^j are plotted in FIG. 2 as functions of $\mathrm{Im}\lambda_5$. The shaded region is prohibited by the vacuum stabilization condition Eq. (5). For the chosen parameters the decays $N_3 \to \sum_a \nu_2 h_a$ give the dominant contribution to ε. Therefore we denote $\varepsilon_3^2 \equiv \varepsilon_{\max}$. As seen in FIG. 2, large CP asymmetries of orded 10^{-3} can be generated *without* resonant condition for heavy neutrinos. The dependence of ε on λ_5 is linear.

In addition, the heavy neutrinos N_i also decay via Eq. (9) with the width $\Gamma_{Di}^{ek} = \Gamma(N_i \to e_{kL}^{\mp}\phi^{\pm})$,

$$\Gamma_{Di}^{ek} = \frac{1}{16\pi} M_{Ni} \epsilon^{2n_{ik}^{\nu}} |y_{ik}^{\nu}|^2 \left(1 - \frac{m_\phi^2}{M_{Ni}^2}\right)^2. \qquad (16)$$

No CP asymmetry is generated in these decays. The total decay rate of N_i is thus given by $\Gamma_{Di} = \Gamma_{Di}^{\nu} + \Gamma_{Di}^{e}$, where $\Gamma_{Di}^{\nu} \equiv \sum_{k=1}^{3} \Gamma_{Di}^{\nu k}$ and $\Gamma_{Di}^{e} \equiv \sum_{k=1}^{3} \Gamma_{Di}^{ek}$.

In order to calculate the generated baryon asymmetry we derive Boltzmann equations for the evolution of heavy neutrino and $\Delta_j = B/3 - L_j$ asymmetry abundances η_{Ni} and η_{Δ_j}, respectively. Here $\eta_X = n_X/n_\gamma$, where n_X is the number density of X and the photon number density is $n_\gamma = 2T^3\zeta(3)/\pi^2$. The Boltzmann equations read

$$\frac{d\eta_{Ni}}{dz} = -\frac{z}{n_\gamma H(z=1)} \left(\frac{\eta_{Ni}}{\eta_{Ni}^{eq}} - 1\right) \gamma_{Di}, \qquad (17)$$

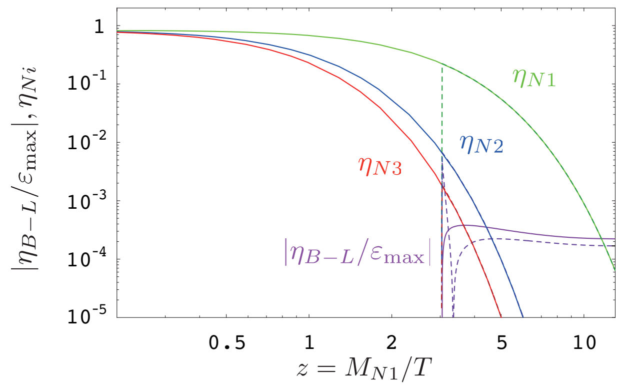

FIG. 3: Evolution of $|\eta_{B-L}/\varepsilon_{\max}|$ and η_{Ni} with $z = M_{N_1}/T$ for $\mathrm{Im}\lambda_5 = 10^{-2}$ in the case of thermal (solid curves) and zero (dashed curves) N_i initial abundances.

$$\begin{aligned} \frac{d\eta_{\Delta j}}{dz} &= -\frac{z}{n_\gamma H(z=1)} \sum_{i=1}^{3} \Bigg[\left(\frac{\eta_{Ni}}{\eta_{Ni}^{eq}} - 1\right) \varepsilon_i^j \gamma_{Di}^{\nu} \\ &- \frac{\eta_{\Delta\nu j}}{2\eta_{\nu_j}^{eq}} \gamma_{Di}^{\nu j} - \frac{\eta_{\Delta ej}}{2\eta_{e_j}^{eq}} \gamma_{Di}^{ej} \Bigg], \end{aligned} \qquad (18)$$

where $\eta_{\Delta j}$ and the asymmetry in left-handed neutrinos (charged leptons) $\eta_{\Delta\nu(e)j}$ are related to each other by the "A-matrix" [10] defined as $\eta_{\Delta\nu(e)j} = A_{jk}^{\nu(e)} \eta_{\Delta k}$ with

$$A^{\nu} = \frac{1}{279} \begin{pmatrix} -89 & 4 & 4 \\ 4 & -89 & 4 \\ 4 & 4 & -89 \end{pmatrix}, \qquad (19)$$

$$A^{e} = \frac{1}{279} \begin{pmatrix} -80 & 13 & 13 \\ 13 & -80 & 13 \\ 13 & 13 & -80 \end{pmatrix}. \qquad (20)$$

The thermally averaged decay rates $\gamma_{Di}, \gamma_{Di}^{\nu(e)}$ and $\gamma_{Di}^{\nu(e)j}$ are obtained from $\Gamma_{Di}, \Gamma_{Di}^{\nu(e)}$ and $\Gamma_{Di}^{\nu(e)j}$ in Eqs.(13) and (16) by $\gamma_{Di}^{\nu(e)j} = n_{Ni}^{\rm eq}(K_1(z_i)/K_2(z_i))\Gamma_{Di}^{\nu(e)j}$, $\gamma_{Di}^{\nu(e)} = \sum_{j=1}^{3} \gamma_{Di}^{\nu(e)j}$ and $\gamma_{Di} = \gamma_{Di}^{\nu} + \gamma_{Di}^{e}$. K_1 and K_2 are the modified Bessel functions and $z_i = M_{Ni}/T$. Hereafter we define $z \equiv z_1$ as usual.

FIG. 3 shows the evolution of $|\eta_{B-L}/\varepsilon_{\max}|$ and η_{Ni} with z for the case of thermal (solid curves) and vanishing (dashed curves) N_i initial abundances. In both cases realistic $B-L$ asymmetry $\eta_{B-L} = \sum_j \eta_{\Delta j}$ is generated below the critical temperature T_c, $z_c \equiv M_{N1}/T_c = 3.04$.

Fast sphaleron processes convert the generated lepton asymmetry to baryon asymmetry [11]. However, sphalerons decay quickly at $T_d < T_c$ [12, 13]. The sphaleron rate $\Gamma_{\Delta(B+L)}$ at temperatures $M_W(T) \ll T \ll M_W(T)/\alpha_W$ is given by $\Gamma_{\Delta(B+L)} \sim M_W(M_W/\alpha_W T)^3 (M_W/T)^3 \exp[-E_{sp}/T]$ [13, 14] where α_W is the $SU(2)_L$ fine structure constant, M_W is W-boson mass and the sphaleron energy is $E_{sp} \sim M_W/\alpha_W$. Just below T_c the sphaleron interactions are faster than the expansion rate of the Universe, $\Gamma_{\Delta(B+L)}/H(z=1) \gg 1$. In this region, η_{B-L} is converted to baryon and lepton asymmetry $\eta_{B,L}$ by sphaleron effect with the

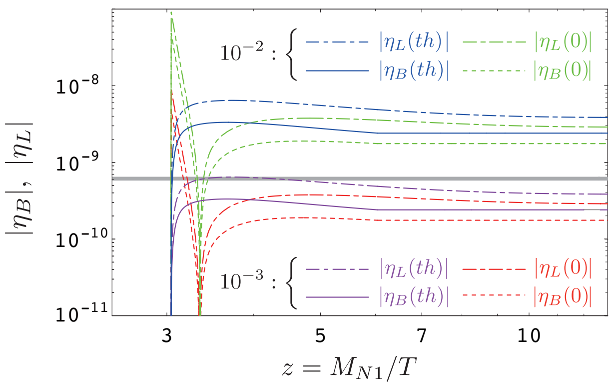

FIG. 4: Lepton and baryon asymmetries for $\mathrm{Im}\lambda_5 = 10^{-2}$ (blue and green curves) and $\mathrm{Im}\lambda_5 = 10^{-3}$ (purple and red curves) for thermal (th) and vanishing (0) initial abundances of N_i at $z = z_c$. The horizontal band is the experimental value $\eta_B^{\rm exp} = (6.05 - 6.37) \times 10^{-10}$ [6].

temperature-depedent rate given by [12, 15]

$$\eta_B = \frac{16T^2 + 10v(T)^2}{46T^2 + 31v(T)^2}\eta_{B-L}, \tag{21}$$

$$\eta_L = -\frac{30T^2 + 21v(T)^2}{46T^2 + 31v(T)^2}\eta_{B-L}. \tag{22}$$

The sphaleron rate $\Gamma_{\Delta(B+L)}$ decreases below T_c by the Boltzmann factor $\exp[-E_{sp}/T]$ and reaches $\Gamma_{\Delta(B+L)}(z_d)/H(z = 1) = 1$ at $z_d = 6.04$. Since the sphaleron processes are effectively switched off at $z > z_d$, the baryon asymmetry is unaffected below z_d.

Numerical results for the evolution of baryon and lepton asymmetries are shown in FIG. 4 for $\mathrm{Im}\lambda_5 = 10^{-2}$ (blue and green curves) and $\mathrm{Im}\lambda_5 = 10^{-3}$ (red and purple curves). Indeed, the sphalerons decouple at z_d while the lepton asymmetry still evolves at $z > z_d$. For appropriate λ_5 the observed baryon asymmetry can be easily generated both for thermal and vanishing initial N_i abundances denoted by (th) and (0), respectively.

Because the leptogenesis window between z_c and z_d is fixed, successful leptogenesis implies an *upper bound* $M_N < 4.5$ TeV on the heavy neutrino mass scale. For larger masses N_i decay above the EWSB scale without being able to generate the observed baryon asymmetry.

CONCLUSIONS

We have presented a novel mechanism for electroweak-scale leptogenesis in the two Higgs doublet scenario of Higgs-dependent Yukawa couplings with global $U(1)_F$ flavor symmetry. In our mechanism the necessary CP-violation comes *entirely* from the Higgs potential, and large CP asymmetries can be *naturally* obtained without resonance enhancement from heavy neutrinos nor any other fine tuning of model parameters. There is a small window of temperatures below T_c where sphalerons are active, yet the observed baryon asymmetry can be easily obtained before sphalerons decay. The requirement of successful leptogenesis implies an *upper bound* $M_N < 4.5$ TeV on the heavy neutrino masses. The distinct prediction [3, 4] of factor of three (five) enhancement of τ, b (μ) Higgs-dependent Yukawa couplings compared to the SM prediction is, in our model, affected only by a small correction related to the generated baryon asymmetry (by λ_5 effects in our numerical example). While discovering TeV scale N_i with small Yukawa couplings at LHC is more than a challenging task [16], this prediction, together with CP-violating Higgs sector testable at LHC [9], makes our scenario of electroweak-scale leptogenesis *realistically* testable at LHC.

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